Flavor Changing Decays of Υ and J/ψ

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We study flavor changing processes $\Upsilon \to B/\overline{B}X_s$ and $J/\psi \to D/\overline{D}X_u$ in the B factories and the Tau-Charm factories. In the standard model, these processes are predicted to be unobservable, so they serve as a probe of the new physics. We first perform a model independent analysis, then examine the predictions of models; such as top color models and MSSM with R-parity violation ; for the branching ratios of $\Upsilon \to B/\overline{B}X_s$ and $J/\psi \to D/\overline{D}X_u$. We find that these branching ratios could be as large as 10^{-6} and 10^{-5} in the presence of new physics.

I. INTRODUCTION

The possibility of observing large CP violating asymmetries in the decay of B mesons has motivated the construction of high luminosity B factories at several of the world's high energy physics laboratories. These B factories will be producing roughly about 10^8 Upsilons. Meanwhile BES has already accumulated 9×10^6 J/ψ and plans to increase the number to 5×10^7 in the near future. An interesting question, that we investigate in this paper, is whether the large sample of the Υ and the J/ψ can be used to probe flavor changing processes in the decays of Υ and J/ψ. In particular we look at the flavor changing processes $\Upsilon \to B/\overline{B}X_s$ and $J/\psi \to D/\overline{D}X_u$, from the underlying $b \to s$ and $c \to u$ quark transitions. For the quarkonium system, these flavor changing processes are expected to be much smaller than in the case of decays of the B or the D meson because of the larger decay widths of the bottonium and the charmonium systems which decay via the strong interactions. Indeed the standard model contributions to $\Upsilon \to B/\overline{B}X_s$ and $J/\psi \to D/\overline{D}X_u$ are tiny. However, new physics may enhance the branching ratios for these processes. Whether this enhancement maybe sufficient for these processes to be observable in the next round of experiments is the subject of this work.

Non leptonic decays of heavy quarkonium systems can be more reliably calculated than the non leptonic decays of the heavy mesons. A consistent and systematic formalism to handle heavy quarkonium decays is available in NRQCD [1] which is missing for the heavy mesons. As in the meson system [2] it is more fruitful to concentrate on quasi-inclusive processes like $\Upsilon \to B/\overline{B}X_s$ and $J/\psi \to D/\overline{D}X_u$ because they can be calculated with less theoretical uncertainty and have larger branching ratios than the purely exclusive quarkonium non leptonic decays. The branching ratios of exclusive flavor changing non leptonic decays of Υ and J/ψ in the standard model have been calculated and found to be very small [3].

We begin with a model independent description of the processes $\Upsilon \to B/\overline{B}X_s$ and $J/\psi \to D/\overline{D}X_u$. In the standard model these decays can proceed through tree and penguin processes. For new physics contribution to these processes we concentrate on four quark operators of the type $\overline{s}b\overline{b}b$ and $\overline{u}c\overline{c}c$. We choose the currents in the four quark operators to be scalars and so these operators may arise through the exchange of a heavy scalar for e.g a Higgs or a leptoquark in some model of new physics. These four quark operators, at the one loop level, generate effective $\overline{s}b\{g, \gamma, Z\}$ and $\overline{u}c\{g, \gamma, Z\}$ vertices which would effect the flavor changing decays of the B and the D mesons. The effective vertices for an on shell g and γ vanish and so there is no contribution to $b \to s\gamma$ or $c \to u\gamma$. We can however put constraints on these operators by considering the processes $b \to sl^+l^-$ and $c \to ul^+l^-$. The constrained operators can then be used to calculate the branching ratios for $\Upsilon \to B/\overline{B}X_s$ and $J/\psi \to D/\overline{D}X_u$.

We then consider some models that may generate the kind of four quark operators described above. A few examples of models where these operators can be generated are top color models and MSSM with R parity violation. In some cases constraints on the parameters that appear in the prediction for the branching ratios for $\Upsilon \to B/\overline{B}X_s$ and $J/\psi \to D/\overline{D}X_u$ are already available. In other cases the parameters are constrained, as in our model independent analysis, from the processes $b \to sl^+l^-$ and $c \to ul^+l^-$.

II. EFFECTIVE HAMILTONIAN.

In the Standard Model (SM) the amplitudes for hadronic Υ decays of the type $b\bar{b} \to s\bar{b} + \bar{s}b$ are generated by the an effective Hamiltonian [4,5] To the standard model contribution we add higher dimensional four quark operators generated by physics beyond the standard model [6]. In this paper, we consider the four quark operators with two scalar currents.

$$L_{new} = \frac{R_1}{\Lambda^2}\bar{s}(1-\gamma^5)b\bar{b}(1+\gamma^5)b + \frac{R_2}{\Lambda^2}\bar{s}(1+\gamma^5)b\bar{b}(1-\gamma^5)b + h.c. \tag{1}$$

The four quark operators in L_{new} are the product of two scalar currents. In Eq. (1) Λ represents the new physics scale and R_1 and R_2 are two free parameters which describe the strength of the contribution of the underlying new physics to the effective operators. In our analysis we will only keep dimension six operators suppressed by $1/\Lambda^2$ and neglect all higher dimension operators. The details of the matrix elements for the processes $\Upsilon \to B/\overline{B}X_s$ and $J/\psi \to D/\overline{D}X_u$ can be found in [6]

III. LOW ENERGY CONSTRAINTS AND MODELS

The lagrangian L_{new} generates, at one loop level, the effective $\bar{s}b\gamma^*$, $\bar{s}bg^*$,$\bar{s}bZ$ vertices, where γ^* and g^* indicate an off shell photon and a gluon. Similar vertices involving $c \to u$ transitions are generated in the charmonium sector also. These vertices, with a γ and Z, will contribute to $b \to sl^+l^-$ and $c \to ul^+l^-$. Note there is no contribution to $b \to s\gamma$. The vertex $b \to sg^*$ can give rise to the process $b \to sq\bar{q}$ which will contribute to non-leptonic B decays. We expect the constraints from $b \to sl^+l^-$ to be better than from non-leptonic B decays because of the theoretical uncertainties in calculating non-leptonic decays The additional contribution to the effective Hamiltonian for $b \to sl^+l^-$ can be written as

$$\delta H_{b\to sl^+l^-} = -\frac{e^2}{16\pi^2}\frac{e_b}{\Lambda^2}\int_0^1 dx 8x(1-x)\log\left(\frac{\Lambda^2}{B^2}\right)\left[R_1\bar{s}\gamma^\mu b_L \bar{l}\gamma_\mu l + R_2\bar{s}\gamma^\mu b_R \bar{l}\gamma_\mu l\right] \tag{2}$$

which has to be added to the standard model contribution [5]. Similar results can also be written for the charm sector.

Now we look at various models that can give rise to L_{new} given in Eq. 1. As a first example we consider a recent version of top color models [7]. In such models the top quark participates in a new strong interaction which is broken at some high energy scale Λ. The strong interaction, though not confining, leads to the formation of a top condensate $< \bar{t}_L t_R >$ resulting in a large dynamical mass for the top quark. The scale Λ is chosen to be of the order of a TeV to avoid naturalness problem which implies that the electroweak symmetry cannot be broken solely by the top condensate. In the low energy sector of the theory, scalar bound states are formed that couple strongly to the b quark [8,9]

$$L_b = \frac{m_t}{f_{\tilde{\pi}}\sqrt{2}}\bar{b}_L(H + iA^0)b_R + h.c \tag{3}$$

where $f_{\tilde{\pi}} \sim 50$ GeV is the top pion decay constant. On integrating out the Higgs fields H and A^0 we have an effective four fermion operator

$$L_{eff} = \frac{m_t^2}{f_{\tilde{\pi}}^2 m_H^2}\bar{b}_L b_R \bar{b}_R b_L \tag{4}$$

Since the b quark in (4) is in the weak-eigenstate, L_{eff} in (4) will induce flavor changing neutral current (FCNC) four quark operators in Eq. (1) after diagonalizing the quark mass matrix [9], with coefficients,

$$\begin{aligned} R_1 &= \frac{1}{4}\frac{m_t^2}{f_{\tilde{\pi}}^2 m_H^2}|D_{Lbb}|^2 D_{Rbb}D^*_{Rbs} \\ R_2 &= \frac{1}{4}\frac{m_t^2}{f_{\tilde{\pi}}^2 m_H^2}|D_{Rbb}|^2 D_{Lbb}D^*_{Lbs} \end{aligned} \tag{5}$$

where D_L and D_R are the mixing matrices in the left and the right handed down sector. In the charm sector similar interactions can arise due to the strong couplings of the top quark to top pions. The effective operators generated by integrating out the top-pions are similar to Eq. (1) with the replacement of b by c and s by u. In topcolor II models [9,10], where there can be strong top-pion couplings of the top with the charm quark, we have

$$
\begin{aligned}
R_1 &= \frac{1}{4}\frac{m_t^2}{f_{\tilde{\pi}}^2 m_{\tilde{\pi}}^2}|U_{Lcc}|^2 U_{Rtc}U_{Rtu}^* \\
R_2 &= \frac{1}{4}\frac{m_t^2}{f_{\tilde{\pi}}^2 m_{\tilde{\pi}}^2}|U_{Rtc}|^2 U_{Lcu}U_{Lcc}^*
\end{aligned}
\tag{6}
$$

In supersymmetric standard models without R parity, the most general superpotential of the MSSM, consistent with $SU(3)\times SU(2)\times U(1)$ gauge symmetry and supersymmetry, can be written as

$$
\mathcal{W} = \mathcal{W}_R + \mathcal{W}_{\not R}, \tag{7}
$$

where $\mathcal{W}_R$ is the R-parity conserving part while $\mathcal{W}_{\not R}$ violates the R-parity. They are given by

$$
\mathcal{W}_R = h_{ij}L_iH_2E_j^c + h'_{ij}Q_iH_2D_j^c + h''_{ij}Q_iH_1U_j^c, \tag{8}
$$

$$
\mathcal{W}_{\not R} = \lambda_{ijk}L_iL_jE_k^c + \lambda'_{ijk}L_iQ_jD_k^c + \lambda''_{ijk}U_i^cD_j^cD_k^c + \mu_iL_iH_2. \tag{9}
$$

Here $L_i(Q_i)$ and $E_i(U_i, D_i)$ are the left-handed lepton (quark) doublet and lepton (quark) singlet chiral superfields, with i, j, k being generation indices and c denoting a charge conjugate field. $H_{1,2}$ are the chiral superfields representing the two Higgs doublets. In the R-parity violating superpotential above, the λ and λ' couplings violate lepton-number conservation, while the λ'' couplings violate baryon-number conservation. λ_{ijk} is antisymmetric in the first two indices and λ''_{ijk} is antisymmetric in the last two indices. While it is theoretically possible to have both baryon-number and lepton-number violating terms in the lagrangian, the non-observation of proton decay imposes very stringent conditions on their simultaneous presence [11]. We, therefore, assume the existence of either L-violating couplings or B-violating couplings, but not the coexistence of both. We calculate the effects of both types of couplings.

In terms of the four-component Dirac notation, the lagrangian involving the λ' and λ'' couplings is given by

$$
\begin{aligned}
\mathcal{L}_{\lambda'} = -\lambda'_{ijk}\Big[&\tilde{\nu}_L^i\bar{d}_R^k d_L^j + \tilde{d}_L^j\bar{d}_R^k\nu_L^i + (\tilde{d}_R^k)^*(\bar{\nu}_L^i)^c d_L^j \\
&-\tilde{e}_L^i\bar{d}_R^k u_L^j - \tilde{u}_L^j\bar{d}_R^k e_L^i - (\tilde{d}_R^k)^*(\bar{e}_L^i)^c u_L^j\Big] + h.c.,
\end{aligned}
\tag{10}
$$

$$
\mathcal{L}_{\lambda''} = -\lambda''_{ijk}\Big[\tilde{d}_R^k(\bar{u}_L^i)^c d_L^j + \tilde{d}_R^j(\bar{d}_L^k)^c u_L^i + \tilde{u}_R^i(\bar{d}_L^j)^c d_L^k\Big] + h.c. \tag{11}
$$

The terms proportional to λ are not relevant to our present discussion and will not be considered here. The exchange of sneutrinos with the λ' coupling will generate L_{new} for $\Upsilon \to \overline{B}X_s$ with

$$
\begin{aligned}
R_1 &= \frac{1}{4}\Sigma_i\frac{\lambda'_{i32}\lambda'^*_{i33}}{m_{\tilde{\nu}_i}^2} \\
R_2 &= \frac{1}{4}\Sigma_i\frac{\lambda'^*_{i23}\lambda'_{i33}}{m_{\tilde{\nu}_i}^2}
\end{aligned}
\tag{12}
$$

For the case of $J/\psi \to DX_u$ the operators in L_{new} cannot be generated at tree level.

IV. RESULTS

The various inputs to our calculations can be found in [6]. The standard model contribution to the branching ratio is 5.2×10^{-11} from the penguin induced $b\to s$ transition. The process $\Upsilon \to \overline{B}X_s$ can also have a contribution in the standard model from tree level processes. The effective Hamiltonian , suppressing the Dirac structure of the currents,

$$H_W = \frac{G_F}{\sqrt{2}} V_{ub} V_{us}^* \left[a_1(\overline{u}b)(\overline{s}u) + a_2(\overline{s}b)(\overline{u}u)\right] \tag{13}$$

where a_1 and a_2 are the QCD coefficients can generate the process $\Upsilon \to B^+K^-$. We can estimate the branching ratio for this process as

$$BR[\Upsilon \to B^+K^-] \approx |\frac{V_{ub}}{V_{cb}}|^2 BR[\Upsilon \to B_c^+K^-]$$

Using $BR[\Upsilon \to B_c^+K^-]$ calculated in Ref [3] one obtains $BR[\Upsilon \to B^+K^-] \sim 1.5 \times 10^{-14}$. For a rough estimate of $BR[\Upsilon \to B^+X_s]$ we can scale $BR[\Upsilon \to B^+K^-]$ by the factor $BR[B \to D^0X]/BR[B \to D\pi]$. The measured value of $BR[B \to D^0X]$ [12] includes D^0 coming from the decay of D^{0*} and D^{+*}. From the spin phase factors $BR[B \to D^*X] \sim 3BR[B \to DX]$. Hence $BR[B \to D^0X]/BR[B \to D\pi] \sim 20$ leading to $BR[\Upsilon \to B^+X_s] \sim 3 \times 10^{-13}$.

So far we have not considered R_1 and R_2 in L_{new}. In our model independent analysis we vary R_1/Λ^2, R_2/Λ^2 one at a time and the use the constraint from measurements of $b \to se^+e^-$ and $b \to s\mu^+\mu^-$. We identify Λ with the masses of the exchange particles which we take to be between $100 - 200$ GeV. The allowed values of R_1/Λ^2, R_2/Λ^2 are then used to calculate $\Upsilon \to \overline{B}X_s$ The constraint from $b \to sl^+l^-$ gives

$$|R_{1,2}|/\Lambda^2 < (6-9) \times 10^{-6}(1/GeV)^2$$

Using the upper bounds on $|R_{1,2}|/\Lambda^2$ we find the branching ratio for the process $\Upsilon(1S) \to \overline{B}X_s$ to be between $(1-2) \times 10^{-6}$. Branching ratios of similar order are also obtained for $\Upsilon(2S)$ and $\Upsilon(3S)$. For $\Upsilon(4S)$ the branching ratio is smaller by a factor of 100 because of the larger width of $\Upsilon(4S)$ which decays predominantly to two B mesons.

Turning now to models, we find for the top color model from Eq. (5) we can write

$$D_{Rbs}^* = 4\frac{R_1}{\Lambda^2}\frac{f_{\tilde{\pi}}^2 m_H^2}{m_t^2 |D_{Lbb}|^2 D_{Rbb}} \tag{14}$$

We can identify $\Lambda = m_H$ and use the constraint from $b \to se^+e^-$ for a typical value of $|R_1|/\Lambda^2 \sim 6 \times 10^{-6}(1/GeV)^2$. Assuming $|D_{Lbb}| \approx |D_{Rbb}| \approx 1$, and $f_{\tilde{\pi}} = 50$GeV we obtain $|D_{Rbs}| \sim 2m_H^2 \times 10^{-6}$. With typical values of $m_H \sim 100 - 200$ GeV we get $|D_{Rbs}| \sim 0.02 - 0.08$. Similar values have been obtained for $|D_{Rbs}|$ in Ref [9] by considering the contributions of the charged higgs and top-pion to $b \to s\gamma$. A similar exercise can be carried out with $|D_{Lbs}|$. Note that B_s mixing probes the combination $D_{Lbs}^* D_{Rbb} D_{Rbs}^* D_{Lbb}$ and so by either choosing $R_1 \sim 0$ or $R_2 \sim 0$ we can satisfy the constraint on $B_s$mixing by choosing the appropriate mixing elements to be small. Note that in top color models we can have operators $\overline{s}(1-\gamma^5)b\overline{d}(1+\gamma^5)d$ and $\overline{s}(1+\gamma^5)b\overline{d}(1-\gamma^5)d$ that can contribute to $\Upsilon \to \overline{B}s\overline{d} \to \overline{B}X_s$ after Fierz reordering. However these operators will be suppressed by form factor effects and also from mixing effects. We have checked that the contribution to $\Upsilon \to \overline{B}X_s$ from these operators are much suppressed relative to the contribution of the operators in L_{new}. We will therefore not consider the the above operators in our analysis.

Turning to R-parity violating susy we first collect the constraints on the relevant couplings. The upper limits of the L-violating couplings for the squark mass of 100 GeV are given by

$$|\lambda'_{kij}| < 0.012, \ (k,j = 1,2,3; i = 1,2), \tag{15}$$
$$|\lambda'_{13j}| < 0.16, \ (j = 1,2), \tag{16}$$
$$|\lambda'_{133}| < 0.001, \tag{17}$$
$$|\lambda'_{23j}| < 0.16, \ (j = 1,2,3), \tag{18}$$
$$|\lambda'_{33j}| < 0.26, \ (j = 1,2,3), \tag{19}$$

The first set of constraints in Eq. (15) come from the decay $K \to \pi\nu\nu$ with FCNC processes in the down quark sector [13]. The set of constraints in Eq. (16) and Eq. (18) are obtained from the semileptonic decays of B-meson [14]. The constraint, on the coupling λ'_{133} in Eq. (17) is obtained from the Majorana mass that the coupling can generate for

the electron type neutrino [15]. The last set of limits in Eq. (19) are derived from the leptonic decay modes of the Z [16]. Assuming all the couplings to be positive we find the branching ratio for $\Upsilon \to \overline{B}X_s$ to be around 2×10^{-6} for $m_{\tilde{\nu}} = 100$GeV.

Turning next to $J/\psi \to \overline{D}X_u$, we first make an estimate for this process in the standard model. Since the penguin $c \to u$ transition is small in the standard model we neglect its contribution. As in the case for the Υ system , for a rough estimate, can write

$$BR[J/\psi \to D^0 X_u] \sim BR[J/\psi \to D^0\pi^0]BR[D^0 \to K^- X]/BR[D^0 \to K^-\pi^+]$$

We obtain $BR[J/\psi \to D^0\pi^0]$ from [3] and keeping in mind that $BR[D^0 \to K^- X]$ contains contributions from states decaying to K^- we obtain $BR[J/\psi \to D^0 X_u] \sim 10^{-10}$. A similar exercise gives $BR[J/\psi \to D^+ X_u] \sim 10^{-9}$.

Considering new physics effects we can constrain R_1 and R_2 from $c \to ul^+l^-$. We get an estimate of the constraint on $c \to ue^+e^-$ by adding up the exclusive modes

$$BR[D \to ue^+e^-] \geq BR[D \to (\pi^0 + \eta + \rho^0 + \omega)e^+e^-]$$

From $c \to ul^+l^-$ one obtains

$$|R_{1,2}|/\Lambda^2 \leq 3.7 \times 10^{-4}(1/GeV)^2$$

We find the branching fraction for the process $J/\psi \to \overline{D}X_u$ using the constraint from $c \to ul^+l^-$ can be $(3-4) \times 10^{-5}$

In top color models taking R_1 and R_2 one at a time, one obtains

$$\frac{2.1 \times 10^3}{m_{\tilde{\pi}}^4}||U_{Lcc}|^2 U_{Rtc}U^*_{Rtu}|^2$$

or

$$\frac{2.1 \times 10^3}{m_{\tilde{\pi}}^4}||U_{Rtc}|^2 U_{Lcc}U^*_{Lcu}|^2$$

as the branching fraction for $J/\psi \to \overline{D}X_u$. For $m_{\tilde{\pi}}$ between $100-200$ GeV this rate can be between $(0.1-2.0) \times 10^{-5}$ if all the mixing angles are ~ 1. It has been shown in Ref [17] that our choices for $f_{\tilde{\pi}}$ and $m_{\tilde{\pi}}$ gives unacceptably large corrections to $Z \to b\overline{b}$ from one loop contribution of the top pions. However in a strongly coupled theory higher loop terms can have significant contributions. Nonetheless if we change $f_{\tilde{\pi}}$ to ~ 100 GeV for better agreement with $Z \to b\overline{b}$ data then the effect in $J/\psi \to \overline{D}X_u$ is reduced by a factor of 16. As in the case of the Υ system we can satisfy the constraint from D mixing by choosing $R_1 \sim 0$ or $R_2 \sim 0$.

For R parity violating susy, contribution to $J/\psi \to \overline{D}X_u$ can only occur at loop level, with both the λ' or λ'' contributing, through the box diagram and so is suppressed.

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